

# Red or blue? A potential kilonova imprint of the delay until black hole formation following a neutron star merger

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Submitted to MNRAS

## ABSTRACT

Mergers of binary neutron stars (NSs) usually result in the formation of a hyper-massive neutron star (HMNS). Whether- and when this remnant collapses to a black hole (BH) depends primarily on the equation of state and on angular momentum transport processes, both of which are uncertain. Here we show that the lifetime of the merger remnant may be directly imprinted in the radioactively powered *kilonova* emission following the merger. We employ axisymmetric, time-dependent hydrodynamic simulations of remnant accretion disks orbiting a HMNS of variable lifetime, and characterize the effect of this delay to BH formation on the disk wind ejecta. Our models follow the system evolution over several seconds, and include the effect of nuclear recombination, viscous heating, and neutrino irradiation by both the HMNS and the disk. When BH formation is relatively prompt ( $\lesssim 100$  ms), outflows from the disk are sufficiently neutron rich to form heavy *r*-process elements with mass number  $A \gtrsim 140$ , resulting in  $\sim$ week-long emission with a spectral peak in the near-infrared (NIR), similar to that produced by the dynamical ejecta. In contrast, delayed BH formation allows neutrinos from the HMNS to raise the electron fraction in the polar direction to values such that potentially *Lanthanide-free* outflows ( $A \lesssim 140$ ) are generated. The lower opacity would produce a brighter, bluer, and shorter-lived  $\sim$  day-long emission (a ‘blue bump’) prior to the late NIR peak from the dynamical ejecta and equatorial wind. A long-lived HMNS also increases the ejecta mass significantly compared to the prompt BH case. Our work motivates efforts to obtain early ( $\lesssim$  day) optical follow-up of mergers detected by Advanced LIGO/Virgo. This new diagnostic of BH formation should be useful for events with a signal to noise lower than that required for direct detection of gravitational waveform signatures.

**Key words:** accretion, accretion disks — dense matter — gravitational waves — hydrodynamics — neutrinos — nuclear reactions, nucleosynthesis, abundances

## 1 INTRODUCTION

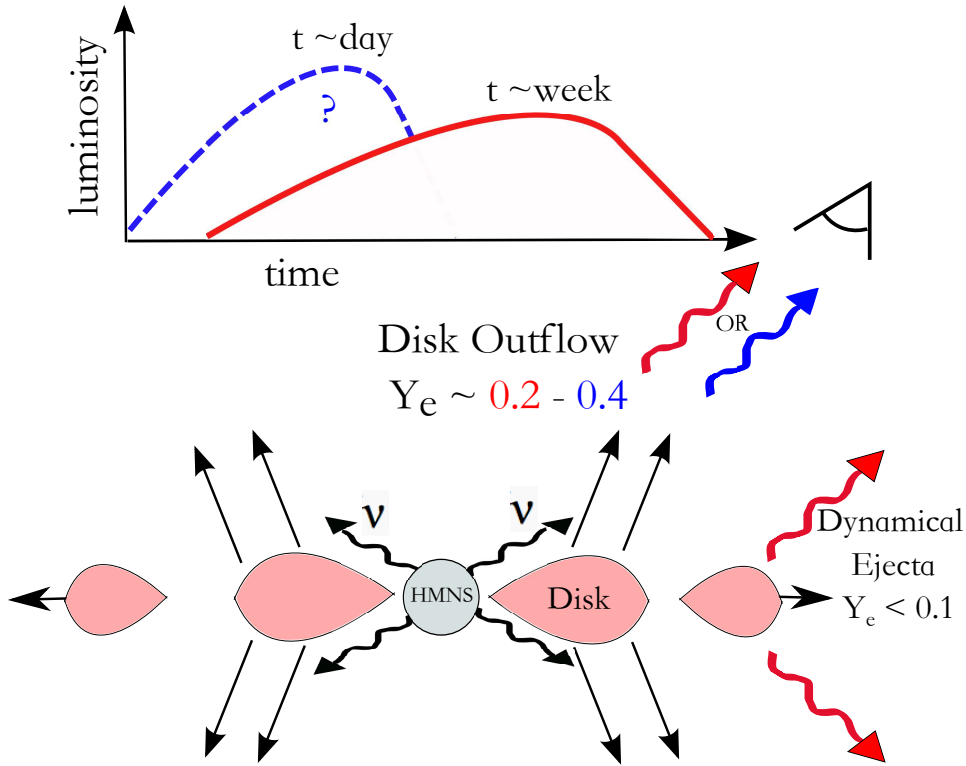
Mergers of binary neutron stars (NSs) (hereafter ‘neutron star mergers’, or NSMs) are the primary source of gravitational waves (GW) for upcoming ground-based interferometric detectors such as Advanced LIGO and Virgo (Abadie et al. 2010). They are also promising central engines for short-duration gamma-ray bursts (GRBs; Paczynski 1986; Eichler et al. 1989; see Berger 2013 for a recent review).

General relativistic simulations of NSMs show that the merger process can result in two qualitatively different outcomes, depending primarily on the total mass of the binary  $M_t$ . If  $M_t$  exceeds a critical value  $M_c$ , then the massive object produced by the merger collapses to a black hole (BH) on the dynamical time ( $\sim$ few ms, e.g., Sekiguchi et al.

2011). On the other hand, if  $M_t < M_c$  then the merger product is at least temporarily supported against gravitational collapse by differential rotation and/or thermal pressure. This meta-stable compact object is usually called a *hypermassive* NS (HMNS; e.g. Kaplan et al. 2013).

The value of  $M_c$  depends on the uncertain equation of state (EoS) of nuclear matter. The recent discovery of massive  $\sim 2M_\odot$  NSs (Demorest et al. 2010; Antoniadis et al. 2013) excludes a soft EoS, placing a lower limit of  $M_c \gtrsim 2.6 - 2.8M_\odot$  (Hotokezaka et al. 2013; Bauswein et al. 2013). It thus appears likely that the ‘canonical’  $1.4 + 1.4 M_\odot$  binary merger goes through a HMNS phase.

When a HMNS does form, its lifetime before collapsing into a BH depends on the timescale for thermal en-



**Figure 1.** Relation between the observed kilonova and the properties of the ejecta that powers it. Material ejected dynamically in the equatorial plane is highly neutron rich ( $Y_e < 0.1$ ), producing heavy  $r$ -process elements that include Lanthanides. This results in emission that peaks in the near-infrared and lasts for  $\sim 1$  week (‘late red bump’) due to the high opacity. Outflows from the remnant disk are more isotropic and also contribute to the kilonova. If the HMNS is long-lived, then neutrino irradiation can increase  $Y_e$  to a high enough value ( $Y_e \sim 0.4$ ) that no Lanthanides are formed, resulting in emission peaking at optical wavelengths (‘early blue bump’). If BH formation is prompt, outflows from the disk remain neutron rich, and their contribution is qualitatively similar to that of the dynamical ejecta.

ergy loss via neutrino emission (e.g., Ruffert & Janka 1999; Paschalidis et al. 2012) and the efficacy of angular momentum transport via gravitational waves and magnetohydrodynamic stresses (e.g., Duez et al. 2006; Stephens et al. 2008; Siegel et al. 2013). For particularly nearby mergers, oscillations excited in the HMNS may be detectable in the GW strain data (Shibata 2005; Bauswein et al. 2012; Hotokezaka et al. 2013). However, the subsequent ring down phase following BH formation is unlikely to be detected by the initial generation of advanced detectors.

Fortunately, NSMs are also accompanied by coincident electromagnetic (EM) signals that inform physical processes at work during the merger (e.g. Metzger & Berger 2012; Kelley et al. 2013; Piran et al. 2013). One such counterpart is a thermal IR/optical transient powered by the radioactive decay of heavy elements synthesized in the merger ejecta (a ‘kilonova’; Li & Paczyński 1998; Metzger et al. 2010; Roberts et al. 2011; Goriely et al. 2011; Piran et al. 2013; Grossman et al. 2013; Tanaka et al. 2014). Kilonovae are particularly promising EM counterparts because (1) their generation is relatively robust, requiring only a modest amount of unbound ejecta; (2) their signal is independent of the existence of a dense surrounding external medium; and (3) unlike a GRB, kilonovae are relatively isotropic. A candidate kilonova was recently detected following the GRB 130603B (Berger et al. 2013; Tanvir et al. 2013).

If the merger ejecta is sufficiently neutron-rich for  $r$ -process nucleosynthesis to reach the Lanthanides ( $A \gtrsim$

139), the optical opacity becomes much higher than that of iron-group elements (Kasen et al. 2013), resulting in emission that is redder, dimmer, and more slowly evolving (Barnes & Kasen 2013; Tanaka & Hotokezaka 2013). Although such unusually red colors may be beneficial in distinguishing NSM transients from unrelated astrophysical sources, the current lack of sensitive wide field infrared telescopes could make EM follow-up across the large sky error regions provided by Advanced LIGO/Virgo even more challenging (e.g. Nissanke et al. 2013; Metzger et al. 2013; Hanna et al. 2013; Kasliwal & Nissanke 2013).

The matter ejected dynamically following a NSM is likely to be sufficiently neutron-rich (as quantified by the electron fraction  $Y_e \lesssim 0.3$ ) to produce a red kilonova (e.g., Rosswog 2005; Duez et al. 2010; Bauswein et al. 2013). Dynamical expulsion is not the only source of ejecta, however. A robust consequence of the merger process is the formation of a remnant torus surrounding the central HMNS. Outflows from this accretion disk over longer, viscous timescales also contribute to the merger ejecta (e.g., Surman et al. 2008; Metzger et al. 2008, 2009a; Lee et al. 2009; Dessart et al. 2009; Wanajo & Janka 2012). The more isotropic geometry of disk winds suggests that they may contribute a distinct component to the kilonova light curve for most viewing angles (Barnes & Kasen 2013; Grossman et al. 2013).

Fernández & Metzger (2013a, hereafter FM13) calculated the viscous evolution of remnant BH accretion disks formed in NSMs using two-dimensional, time-dependent hy-

hydrodynamical simulations. Over several viscous times, FM13 found that a fraction  $\sim$ several percent of the initial disk mass is ejected as a moderately neutron-rich wind ( $Y_e \sim 0.2$ ) powered by viscous heating and nuclear recombination. Although the higher entropy of the outflow as compared to the dynamical ejecta results in subtle differences in composition (e.g. a small quantity of helium), the disk outflows likely produce Lanthanide elements with sufficient abundance to result in a similarly red kilonova as with the dynamical ejecta.

FM13 included the effects of self-irradiation by neutrinos on the dynamics and composition of the disk. Due to the relatively low accretion rate and radiative efficiency at the time of the peak outflow, neutrino absorption had a sub-dominant contribution to the disk evolution. This hierarchy is important because a large neutrino flux tends to drive  $Y_e$  to a value higher than that in the disk midplane (e.g. Surman et al. 2008; Metzger et al. 2008; Surman et al. 2013). If neutrino irradiation is sufficient to drive  $Y_e \gtrsim 0.3 - 0.4$ , the nuclear composition of the disk outflows would be significantly altered, resulting in a distinct additional component visible in the kilonova emission.

By ignoring the influence of a central HMNS, FM13 implicitly assumed a scenario in which BH formation was prompt or the HMNS lifetime very short. Here we extend the study of FM13 to include the effects of neutrino irradiation from a long-lived HMNS. As we will show, the much larger neutrino luminosity of the HMNS has a profound effect on the quantity and composition of the disk outflows, allowing a direct imprint of the HMNS lifetime on the kilonova (Figure 1). As in FM13, our study includes many approximations that enable us to follow the secular evolution of the system. We focus here on exploring the main differences introduced by the presence of a HMNS, and leave more extensive parameter space studies or realistic computations for future work.

The paper is organized as follows. In §2 we describe the numerical model employed. Our results are presented in §3, separated into dynamics of the outflow (§3.1) and composition (§3.2). A summary and discussion follows in §4. Appendix A describes in more detail the upgrades to the neutrino physics implementation relative to that of FM13.

## 2 NUMERICAL MODEL

Our numerical model largely follows that described in FM13. Here we summarize the essential modifications needed to model the presence of a HMNS.

### 2.1 Equations and Numerical Method

We use FLASH3.2 (Dubey et al. 2009) to solve the time-dependent hydrodynamic equations in two-dimensional, axisymmetric spherical geometry. Source terms include the pseudo-Newtonian potential of Paczyński & Wiita (1980), an anomalous shear stress for angular momentum transport (§2.2), and charged-current weak interaction terms in the energy and lepton number equations (§2.3). The equation of state is that of Timmes & Swesty (2000), with the abundances of neutrons, protons, and alpha particles satisfying nuclear statistical equilibrium, and including the nuclear

binding energy of alpha particles in the internal energy<sup>1</sup>. The upper density limit in the tabulated lepton component is extended using analytic expressions for a fully relativistic, arbitrary-degeneracy lepton gas (e.g., Bethe et al. 1980).

The code contains modifications relative to the public version, aimed at investigating the viscous evolution of merger remnant accretion disks (Fernández 2012; Fernández & Metzger 2013b; FM13). The radial cell spacing is non-uniform, with consecutive cells having a constant ratio between their sizes. The meridional grid is uniform in  $\cos\theta$ . Models are initialized from an equilibrium torus that is allowed to relax for 100 orbits – without source terms – to smooth out initial discontinuities.

### 2.2 Boundary Conditions and Angular Momentum Transport

We approximate the HMNS by a reflecting inner boundary at a fixed radius  $R_{\text{NS}} = 30$  km for all variables except specific angular momentum, motivated by the approximate location of the neutrosphere in the models of Dessart et al. (2009). Our approximation is justified in that HMNSs from self-consistent calculations achieve a state of quasi-equilibrium over a few dynamical times after the merger, particularly if they are not too close to the mass for prompt BH formation (e.g., Sekiguchi et al. 2011). While our spherical boundary does not initially follow the highly elliptical form of isobaric surfaces, the accumulation of matter above the boundary develops such an elliptical shape after a few orbits of evolution. Whenever the NS is assumed to collapse to a BH, this reflecting boundary condition is changed to absorbing for all variables as in FM13. The outer radial boundary allows matter to leave the domain, while the symmetry axis is reflecting in the meridional direction.

Angular momentum transport is mediated by an anomalous shear stress as in FM13. The coefficient of kinematic viscosity is that of Shakura & Sunyaev (1973). Whenever the HMNS is present, the specific angular momentum in the inner radial ghost cells is set so that uniform rotation with an angular velocity  $\Omega_* = 2\pi/P_*$  is obtained (here  $P_*$  is the assumed rotation period of the neutron star). The viscous stress is applied at the inner radial boundary, thus driving the inner most active cells to co-rotate with the neutron star, and resulting in the formation of a boundary layer (e.g., Frank et al. 2002).<sup>2</sup> Although the MRI is unlikely to operate in the boundary layer itself, waves produced by this interface can result in angular momentum transport of a similar magnitude (Belyaev et al. 2013).

The steep density gradient that develops in the boundary layer requires verifying that results are converged. We elaborate on this in Appendix A.

<sup>1</sup> We have corrected an error in the treatment of the nuclear binding energy of  $\alpha$  particles, which led to an overestimation of the recombination heating in the models of FM13. With the correct treatment, the amount of mass ejected decreases by a factor of  $\sim$ few and the ejection occurs over a longer period of time relative to the values published in FM13. Changes to other properties relevant for nucleosynthesis ( $Y_e$ , entropy) are insignificant.

<sup>2</sup> We neglect the spin-up (or spin-down) of the neutron star by this torque, however, because the angular momentum contained in the disk is small compared to that in the star.

**Table 1.** Models Evolved and Outflow Properties. The first six columns show model name, HMNS lifetime, stellar period, fractional radial resolution at inner boundary, and irradiation by the neutron star, respectively. The last eight columns show integrated properties of the outflow, restricted to equatorial ( $\theta \in [30^\circ, 150^\circ]$ ) and polar ( $\theta < 30^\circ$  and  $\theta > 150^\circ$ ) latitudes. Each group of four columns includes the mass-flux weighted electron fraction, entropy, and expansion time at a radius where the mass-flux weighted temperature is  $5 \times 10^9$  K, as well as the net mass in unbound material crossing a surface at  $10^9$  cm, normalized by the initial torus mass.

Model	$t_{\text{ns}}$ (ms)	$P_\star$	$\alpha$	$\Delta r_{\text{min}}/R_\star$	$\star\text{-Irr.}$	Equatorial Outflow				Polar Outflow <sup>b</sup>			
						$\bar{Y}_e$	$\bar{s}$ (k/b)	$t_{\text{exp}}$ (ms)	$M_{\text{ej,unb}}/M_{\text{t0}}$ (%)	$\bar{Y}_e$	$\bar{s}$ (k/b)	$t_{\text{exp}}$ (ms)	$M_{\text{ej,unb}}/M_{\text{t0}}$ (%)
t000A3	0	...	0.03	5E-3	No	0.18	18	27	2.1	...	...	...	0.9
t010A3p15	10	1.5			Yes	0.19	20	30	2.5	...	...	...	1.0
t030A3p15	30					0.23	21	27	3.9	0.50	45	7.6	1.3
t100A3p15	100					0.28	20	33	9.2	0.47	42	11	4.0
t300A3p15	300					0.30	20	32	28	0.44	38	17	8.8
tInfA3p15	$\infty$					0.32	21	42	67	0.43	52	17	22
pA1p15	$\infty$	1.5	0.01	5E-3	Yes	0.36	21	61	47	0.45	52	19	22
pA3p20		2	0.03			0.30	29	37	62	0.37	62	15	17
pA3p15xs		1.5			No	0.27	21	45	56	0.35	61	15	23
rA3p15r1	$\infty$	1.5	0.03	1E-2	Yes	0.32	21	45	69	0.43	53	17	22
rA3p15r2 <sup>a</sup>				2E-3		...	...	...	...	...	...	...	...

<sup>a</sup> Model rA3p15r2 was evolved for a shorter time. See Appendix A for details.

<sup>b</sup> Average thermodynamic quantities are not computed if there is no region where the average temperature is  $\sim 5 \times 10^9$  K.

### 2.3 Neutrino Treatment

We consider contributions to the neutrino irradiation from both the central HMNS and from the disk. The neutrino flux from the HMNS is assumed to be spherically symmetric and isotropic, following a Fermi-Dirac spectrum with zero chemical potential and a constant neutrinospheric temperature. We take equal luminosities of electron neutrinos and antineutrinos, with temperatures  $T_{\nu_e, \text{ns}} = 4$  MeV and  $T_{\bar{\nu}_e, \text{ns}} = 5$ , respectively, yielding mean neutrino energies of  $\sim 12$  and  $15$  MeV, respectively. We ignore the highly aspherical form of the neutrino flux and temperature obtained in simulations with more sophisticated neutrino transport (e.g., Dessart et al. 2009). Since the viscously-driven winds arise on timescales of a few seconds, we must account for the change in the neutrino luminosity due to the cooling evolution of the HMNS (e.g., Roberts 2012). We do this by adopting a temporal fit to the cooling curves of Pons et al. (1999), with a normalization that approximately matches the luminosities at 30 ms obtained by Dessart et al. (2009),  $2 \times 10^{52}$  erg s<sup>-1</sup>. The neutrino flux incident upon a given point in the computational domain is attenuated by the integrated optical depth from the inner boundary along radial rays.

Given the large optical depths obtained in the boundary layer, we replace the neutrino cooling implementation used in FM13 – which was appropriate for a disk of low to moderate optical depth – with a neutrino leakage scheme that follows that of Ruffert et al. (1996). The spatial emission properties of the disk itself are also slightly more complicated when a HMNS sits at the center than when a black hole is the central object, thus our prescription for disk *self-irradiation* must be modified. A detailed description of these changes is documented in Appendix A.

Charged-current weak interaction rates are otherwise unmodified relative to FM13. Additional neutrino emission channels are sub-dominant and hence are neglected. We impose a floor on the electron fraction at  $Y_{e, \text{min}} = 0.01$  to

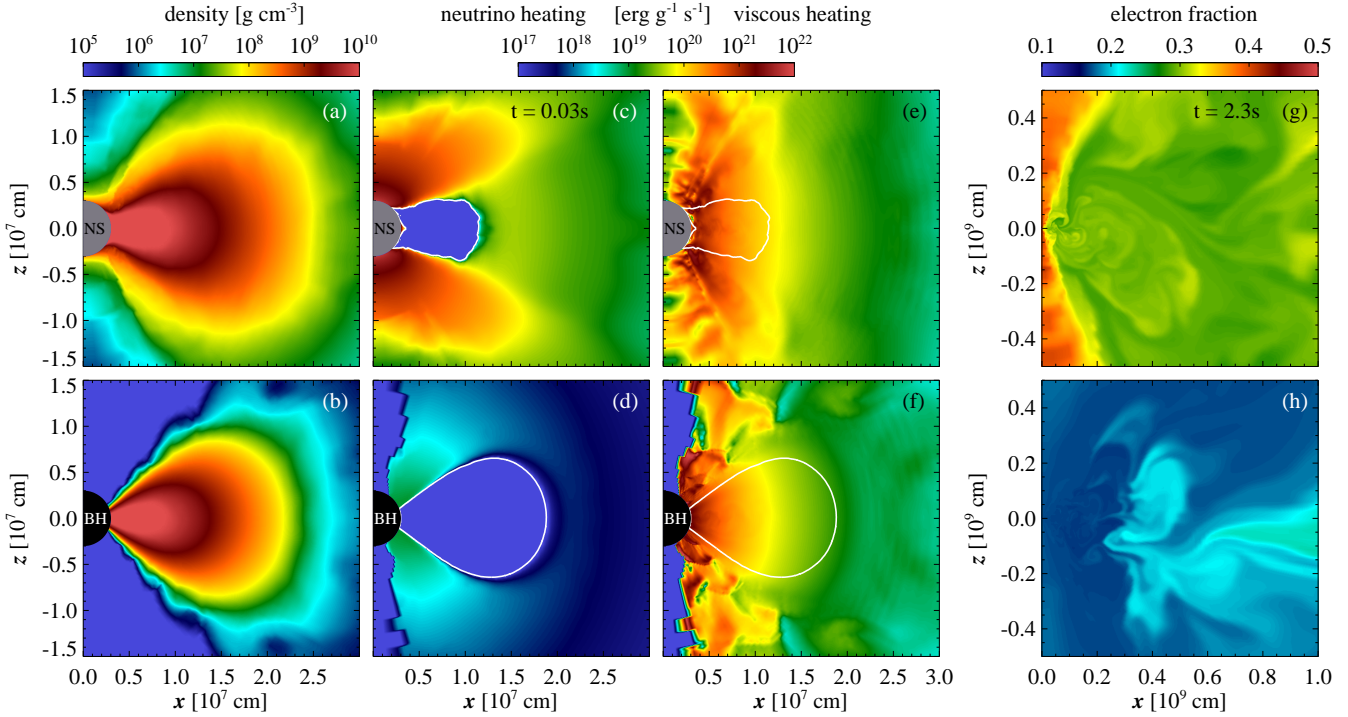
prevent problems with our tabulated rates. Finally, we neglect the neutrino contributions to the pressure and internal energy in optically thick regions: this contribution is at most  $\sim 10\%$  in the equatorial part of the boundary layer.

### 2.4 Models Evolved

The models run are summarized in Table 1. Our baseline parameter set consists of a central object (BH or HMNS) with mass  $3M_\odot$ , a disk mass  $M_{\text{t0}} = 0.03M_\odot$ , constant specific angular momentum, and a ‘disk radius’  $R_0 = 50$  km. The entropy and electron fraction are uniform, initially set to  $8k_B$  per baryon and 0.1, respectively. We use 384 cells in radius, with an innermost radial spacing set to  $\delta r_{\text{min}}/R_\star \simeq 5 \times 10^{-3}$  in the fiducial set of models. A total of 112 cells are used in the  $\theta$  direction. Most models are evolved for 3000 orbits at  $r = R_0$ , or  $\sim 8.7$  s.

Six models probe the effect of the HMNS lifetime  $t_{\text{ns}}$  on the mechanical and compositional properties of the disk wind (t-series). The model with zero HMNS lifetime (t000A3, pure BH) has the same set of parameters as the fiducial model used in FM13, but differs in its neutrino implementation (§2.3) and in the treatment of the nuclear binding energy of  $\alpha$  particles (§2.1).

Another set of three models (p-series) probes the effect of varying important parameters relative to the baseline model with an infinite HMNS lifetime (tInfA3p15). We explore the sensitivity to the assumed rotation period  $P_\star$ , the magnitude of the viscous stress  $\alpha$ , and the effect of removing neutrino irradiation by the HMNS. Finally, two models (r-series) quantify the sensitivity of our results on the radial resolution employed near the inner boundary (Appendix A).



**Figure 2.** Snapshots in the evolution of models tInfA3p15 (stable HMNS, top) and t000A3 (prompt BH, bottom). The nine contiguous panels on the left show density (a,b), neutrino heating (c,d), and viscous heating (e,f) at time 0.03s, corresponding approximately to 10 orbits at the initial density peak. The two panels on the right (g,h) show the electron fraction at time 2.3s (800 orbits at the initial density peak). The white contours show the gain surface, inside which emission of neutrinos and antineutrinos dominates over their absorption by nuclei. Models suppress neutrino and viscous source terms below a density of  $10 \text{ g cm}^{-3}$  (FM13) for numerical reasons (as in the polar region of panels d and f).

### 3 RESULTS

#### 3.1 Evolution and Mass Ejection

The presence of a HMNS results in two important differences relative to the case where a BH forms promptly. To illustrate these differences, we first focus the discussion on models tInfA3p15 and t000A3, corresponding to infinitely lived HMNS and prompt BH, respectively.

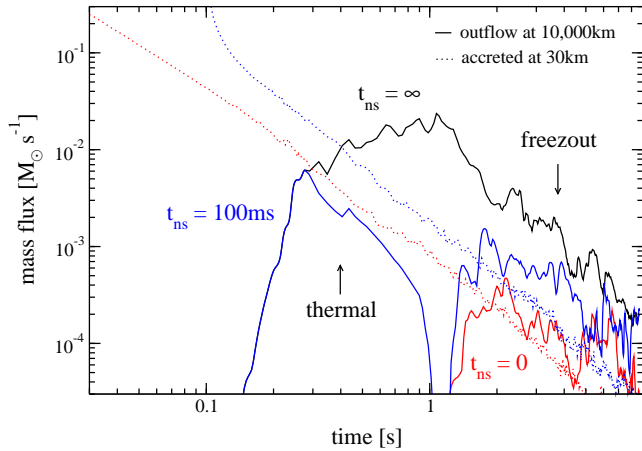
First, a hard stellar surface leads to the formation of a boundary layer over the first few orbits as the disk begins to accrete. Material that has angular momentum removed by the viscous stress migrates towards the polar regions of the stellar surface, forming an envelope (Figure 2a). In contrast, a BH accretion disk is such that the density undergoes a steep decline at an angle from the midplane, with the polar regions being largely devoid of significant material (Figure 2b).

The second important difference introduced by the HMNS is the level of neutrino heating in the system. The geometry of this heating is shown in Figure 2c. Irradiation by the HMNS is strongest in the polar regions, where there is less attenuating material. Emission in this region is also contributed to by the disk itself, which emits mostly from two emission spots above the midplane near the stellar surface. While regions near the disk midplane are dominated by neutrino cooling at inner radii, neutrino heating dominates cooling outside  $\sim 100 \text{ km}$  even when the irradiation from the star is shadowed by the disk. These emission properties stand in stark contrast to the case of a BH (Figure 2d), in

which neutrino heating is largely confined to the polar regions, and at a much lower level, which makes it insignificant relative to viscous heating (FM13).

The combined action of these two effects – boundary layer and higher degree of neutrino heating – leads to very different mass ejection properties depending on the nature of the central object (Figure 3). On the one hand, the mass ejection mechanism from a BH accretion disk relies on the *weak freezeout* of the disk (e.g., Metzger et al. 2009b), in which neutrino cooling shuts down as the disk spreads viscously and the temperature decreases. Viscous heating, nuclear recombination, and transport of rotational kinetic energy outwards all act as positive energy source terms which are uncompensated at radii  $\gtrsim 100 \text{ km}$ , leading to ejection of the outer layers of the disk. This ejection manifests as a broad equatorial outflow ( $\sim 60^\circ$  from the equator) that carries several percent the initial disk mass on a timescale of a few seconds (FM13).

In contrast, a HMNS disk displays a distinct phase of mass ejection that operates on the thermal time of the outer disk ( $\sim 30 - 100 \text{ ms}$  for  $r \sim 100 - 200 \text{ km}$ ) due to the larger amount of neutrino heating and the concentration of neutrino cooling in a smaller spatial region (Fig. 2c). This phase occurs well before weak freezeout. The presence of a reflecting boundary condition also implies that inward traveling sound waves – generated when material is heated and expands – are available for pressure buildup instead of being absorbed by a BH. The *equatorial outflow* is thus signifi-



**Figure 3.** Mass loss rate in unbound material at 10,000 km (solid lines) and net accretion rate at 30 km (dotted lines) as a function of time. Shown are three models that illustrate the difference between prompt BH formation (t000A3p15, red), infinitely lived HMNS (tInfA3p15, black), and an intermediate case with  $t_{\text{ns}} = 100$  ms (t100A3p15). Mass fluxes are computed over the full range of polar angles. Mass ejection before 1 s occurs on the thermal time of the disk, while the late-time wind happens due to weak freezout (Metzger et al. 2009b).

cantly enhanced when the HMNS is present for a time much longer than the thermal time of the disk.

In addition, strong neutrino heating at polar latitudes in HMNS disks causes a genuine *neutrino-driven wind* at early times in this direction (see, e.g., Dessart et al. 2009). After several 10 ms, viscous heating takes over neutrino heating. Overall, mass loss along the polar directions goes from less than 1% of the disk mass in the pure BH case to nearly  $\sim 20\%$  for an infinitely lived HMNS (Table 1). Given that this material is strongly irradiated, its composition is qualitatively different from that ejected along the disk midplane (§3.2).

Allowing the HMNS to survive for different lengths of time results in a behavior that is intermediate between the pure BH and infinitely lived HMNS, as shown in Figure 3 for a HMNS with a lifetime of 100 ms. The ejected mass depends monotonically on the HMNS lifetime (Table 1), at both polar and equatorial latitudes. Polar outflows are significant if  $t_{\text{ns}} \gtrsim 100$  ms.

The dynamics of the flow at the transition from HMNS to BH are characterized by instantaneous accretion of the boundary layer material, and subsequent generation of a rarefaction wave. The velocity becomes negative from inside out, and the initial phase of accretion is interrupted. This cutoff in the mass loss is quite steep, as shown in Figure 3 for the model with  $t_{\text{ns}} = 100$  ms. The equatorial regions of the disk relax to the BH configuration, and the late-time outflow appears at the expected time of  $\sim 1$  s. The fluid at polar latitudes, which was previously supported mostly by gas pressure, collapses into the hole leaving an evacuated polar cavity. A fraction of the material still manages to escape, however, as shown in Table 1.

Decreasing the viscosity of the disk (model pA1p15) relative to the fiducial  $t_{\text{ns}} = \infty$  case (tInfA3p15) results in a smaller amount of equatorially ejected material. The main reason for this difference lies in that viscous heating is more

centrally concentrated for lower  $\alpha$ . The accretion luminosity is also smaller by a factor of  $\sim 2$ , resulting in less neutrino heating in equatorial regions on the thermal timescale of the disk. The time-integrated neutrino energy deposition at polar latitudes differs only by  $\sim 10\%$ , however.

Increasing the rotation period of the star (model pA3p20) results in less mass ejected along the poles, with a smaller decrease in the equatorial ejection relative to the fiducial model. This result can be traced back to the behavior of the boundary layer. Since the fluid is forced to co-rotate with the star at the surface, a lower rotation rate implies a lower centrifugal force, which allows the boundary layer to spread more material to high latitudes (the component of the centrifugal force tangential to the stellar surface points towards the midplane; e.g., Inogamov & Sunyaev 2010). This excess of material is able to attenuate more neutrino flux, decreases the heating, resulting in less mass ejection (with a smaller electron fraction).

Finally, removing irradiation from the HMNS while keeping the reflecting boundary and self-irradiation from the disk (model pA3p15xs) results in a lower amount of mass ejected in equatorial regions. As in the case of lower  $\alpha$ , there is less neutrino heating at larger radii on the disk midplane. The decrease in the equatorial mass ejection relative to the fiducial model is less than in the low- $\alpha$  case, however, because viscous heating is unchanged. The fact that the amount of mass ejected along the poles is comparable (with smaller  $Y_e$ , however) shows that viscous heating is also a fundamental agent in driving the polar outflow.

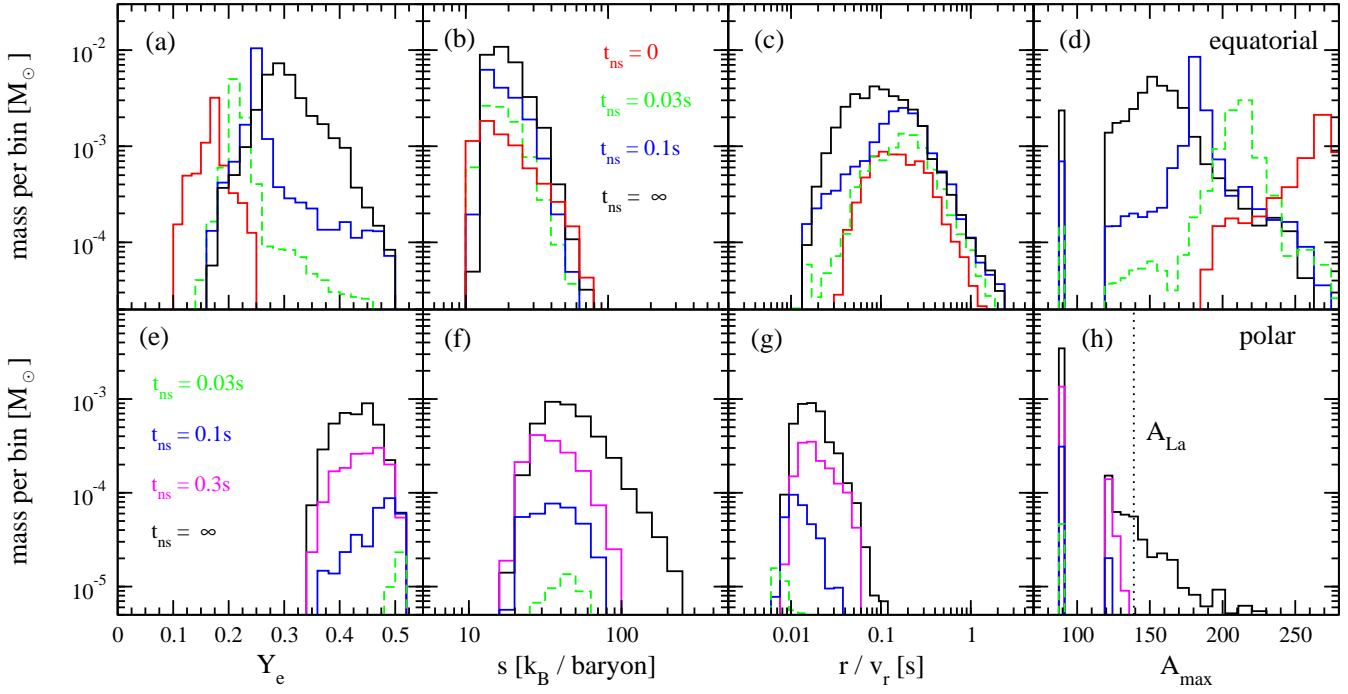
### 3.2 Composition

Free nuclei recombine into  $\alpha$ -particles once the temperature decreases to  $T \lesssim 10^{10}$  K. Heavier elements start to form once the temperature decreases further,  $T \lesssim 5 \times 10^9$ , via the reaction  ${}^4\text{He}(\alpha, \gamma){}^8\text{Be}(\alpha, n){}^{12}\text{C}$ . After  ${}^{12}\text{C}$  forms, additional  $\alpha$ -captures produce heavy ‘seed’ nuclei with characteristic mass  $\bar{A} \simeq 90 - 120$  and charge  $\bar{Z} \simeq 35$  (the ‘ $\alpha$ -process’; Woosley & Hoffman 1992). Whether nucleosynthesis proceeds to heavier  $r$ -process nuclei of mass  $A$  depends on the ratio of free neutrons to seed nuclei once the  $\alpha$ -process completes. Since the formation of  ${}^{12}\text{C}$  is the rate-limiting step in forming seeds, this critical ratio depends primarily on three quantities (e.g. Hoffman et al. 1997; Appendix C of FM13): the electron fraction  $Y_e$ , entropy  $S$ , and the expansion timescale  $t_{\text{exp}}$  at times just following  $\alpha$  particle formation ( $T \sim 5 \times 10^9$  K).

A mass-flux average of each of these three quantities is given in Table 1 for most models (see FM13 for a description of the calculation method). The radial position for the average is chosen so that the mass-flux averaged temperature is approximately  $5 \times 10^9$  K. The average is separated between equatorial and polar latitudes ( $60^\circ$  from the midplane and  $30^\circ$  from the axis, respectively). In the case of prompt BH formation or very shortly lived HMNS ( $t_{\text{ns}} \leq 10$  ms), there is never enough material in the polar regions to achieve the desired temperature, and hence the average is not computed there. After collapse to a BH, the polar region is evacuated, so the radius for the average is obtained for times less than the HMNS lifetime (but averaged thermodynamic quantities are computed using the whole evolution).

The average electron fraction of the material is a mono-





**Figure 4.** Mass histograms of material with positive velocity, evaluated at radii such that the mass-flux weighted temperature is  $\sim 5 \times 10^9$  K. The counting is restricted to latitudes within  $60^\circ$  of the midplane (top) and within  $30^\circ$  of the polar axis (bottom). Shown are the distributions of electron fraction (a,e), entropy (b,f), expansion time (c,g), and maximum mass number of  $r$ -process elements (d,h) obtained from equations (3)-(6) assuming seed nuclei with  $\bar{A} = 90$  and  $\bar{Z} = 36$ . Different curves correspond to models with different HMNS lifetime, as labeled. The polar outflow has a peak  $A_{\max}$  below the Lanthanides ( $A \gtrsim A_{\text{La}} = 139$ ). The gap above the seed mass number arises from the steep drop of  $X_n$  with  $Y_e > \bar{Z}/\bar{A}$  at these entropies (eq. 4).

tonic function of the HMNS lifetime. This is a direct consequence of the higher level of neutrino irradiation introduced by the HMNS. A rough estimate of the change in  $Y_e$  over a thermal time in the boundary layer yields

$$\begin{aligned} \Delta Y_e &\sim \frac{Q_\nu \Delta t}{\langle \varepsilon_\nu \rangle / m_n} \\ &\sim 1 \left( \frac{Q_\nu}{10^{21} \text{ erg g}^{-1} \text{ s}^{-1}} \right) \left( \frac{10 \text{ MeV}}{\langle \varepsilon_\nu \rangle} \right) \left( \frac{\Delta t}{10 \text{ ms}} \right), \end{aligned} \quad (1)$$

where  $Q_\nu$  is the specific neutrino heating rate (Figure 2c),  $\langle \varepsilon_\nu \rangle$  is the mean neutrino energy, and  $\Delta t$  is the time interval. Irradiation by the HMNS can thus introduce changes of order unity in the electron fraction over the time it takes the polar outflow to be launched. Similar considerations apply to the equatorial outflow, although the fact that this ejecta originates in regions of the disk that are both farther out in radius and shadowed by the inner regions causes the changes in  $Y_e$  to be less pronounced.

It is worth keeping in mind that we are imposing equal luminosities of electron neutrinos and antineutrinos from the HMNS (§2.3). In the limit of high irradiation, neutrinos drive  $Y_e$  towards (Qian & Woosley 1996)

$$Y_e^{\text{eq}} \simeq \left[ 1 + \frac{\epsilon_{\bar{\nu}_e} - \Delta + \Delta^2 / \epsilon_{\bar{\nu}_e}}{\epsilon_{\nu_e} + \Delta + \Delta^2 / \epsilon_{\nu_e}} \right]^{-1}, \quad (2)$$

where  $\Delta = 1.293$  MeV is the neutron-proton mass difference, and  $\epsilon_{\nu_i}$  is the ratio of the mean square energy to the mean energy of the distribution,  $\langle E_\nu^2 \rangle / \langle E_\nu \rangle$ . Our Fermi-Dirac spectrum with zero chemical potential implies  $\epsilon_{\nu_i} \simeq 4kT_{\nu_i}$ , hence  $Y_e^{\text{eq}} \simeq 0.52$ .

The distribution of thermodynamic properties of material with positive velocity – evaluated at the averages in Table 1 – is shown in Figure 4. The electron fraction of the equatorial material has a peak that tracks the average value, with a long tail to high  $Y_e$  for longer HMNS lifetime. In contrast, the polar material has more material with lower  $Y_e$  for longer  $t_{\text{ns}}$ . This is a reflection of the time-dependence of the outflow composition, which is dominated by neutron-rich conditions at late time (Figure 5).

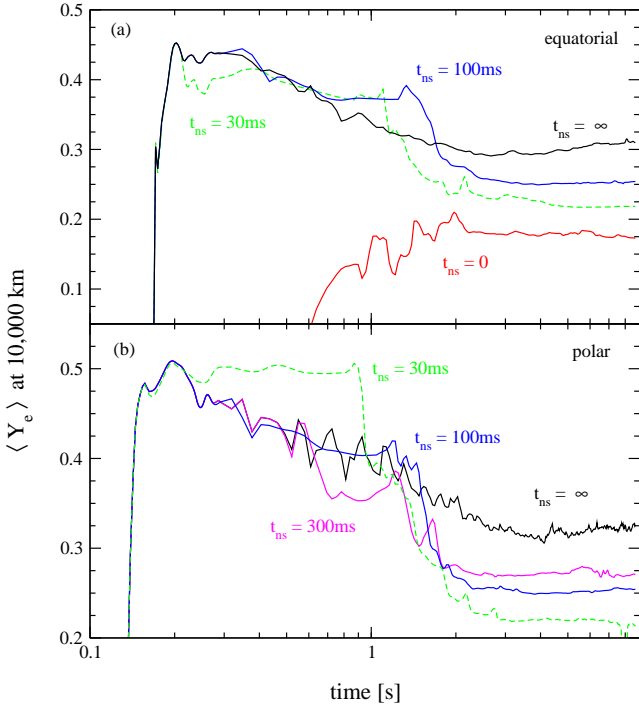
The distribution of entropy in equatorial material is insensitive to the lifetime of the HMNS. The polar material also has its peak nearly unchanged, though a high-entropy tail develops for longer  $t_{\text{ns}}$ . A similar behavior is seen in the expansion time, with a peak values that undergo moderate shifts for both polar and equatorial outflows, but with changes in the shape of the distribution for longer-lived HMNSs.

The maximum mass number  $A_{\max}$  to which the  $r$ -process may proceed can be estimated as (Hoffman et al. 1997)

$$A_{\max} = \bar{A} \left[ 1 + \frac{X_n}{1 - X_n - X_{\text{He}}} \right], \quad (3)$$

where  $X_n$  and  $X_{\text{He}}$  are the mass fractions of neutrons and alpha particles, respectively, following the  $\alpha$ -process. These mass fractions can be estimated analytically in two limits, depending on whether  $Y_e$  is greater or less than the electron fraction of the seed nuclei  $\bar{Y}_e = \bar{Z}/\bar{A}$  (Hoffman et al. 1997):

$$X_n \simeq \begin{cases} 1 - 2Y_e, & (Y_e < \bar{Y}_e) \\ (1 - 2Y_e) \exp[-(\bar{A} - 2\bar{Z})(Y_e/2)^3 \mathcal{F}] & (Y_e > \bar{Y}_e), \end{cases} \quad (4)$$



**Figure 5.** Mass-flux weighted electron fraction at 10,000 km as a function of time, for models with different HMNS lifetime  $t_{\text{ns}}$ . The averages are restricted to latitudes within  $60^\circ$  of the midplane (top) and within  $30^\circ$  of the polar axis (bottom). The drop around 1 – 2 s for the polar material arises from the combination of a decreasing neutrino luminosity and the travel time to 10,000 km. Note also that the vertical scales are different.

$$X_{\text{He}} \simeq \begin{cases} \left[ \frac{\bar{Z}(1 - 2Y_e)\mathcal{F}/16 + (2Y_e)^{-2}}{2Y_e} \right]^{-1/2} & (Y_e < \bar{Y}_e) \\ Y_e & (Y_e > \bar{Y}_e), \end{cases} \quad (5)$$

where the factor

$$\mathcal{F} = 8000 \left( \frac{t_{\text{exp}}}{100\text{ms}} \right) \left( \frac{s}{20k_{\text{B}} \text{ nuc}^{-1}} \right)^{-3} \quad (6)$$

quantifies the yield of the reaction  ${}^4\text{He}(\alpha, \gamma){}^9\text{Be}(\alpha, n){}^{12}\text{C}$  in the expanding outflow.

The resulting distribution of  $A_{\text{max}}$  (eq. [3]) for equatorial and polar outflows, obtained assuming  $\bar{A} = 90$  and  $\bar{Z} = 35$ , is shown<sup>3</sup> in Figure 4. The equatorial wind is expected to make heavy r-process elements, with the peak  $A_{\text{max}}$  decreasing with longer HMNS lifetime down to  $\sim 150$  for  $t_{\text{ns}} \rightarrow \infty$ . The distribution is broad, however, with a characteristic width  $\Delta A_{\text{max}} \sim 50$ . Note that a value of  $A_{\text{max}} \gg 200$  does not necessarily imply that such heavy nuclei will form, but rather that some fraction will fission to lighter nuclei  $A \sim 130$  before capturing additional neutrons. The gap from the seed mass number to  $A \simeq 120$  arises due to the steep drop of  $X_n$  with  $Y_e > \bar{Y}_e$  at these entropies.<sup>4</sup>

In contrast, the polar outflow maintains the composi-

tion of the seed nuclei, with a minor contribution at higher mass number in the case of a long-lived HMNS. As shown in Figure 5, in all cases the late-time value of  $Y_e$  is lower than that in the initial episode of mass ejection. Hence the tail in  $A_{\text{max}}$  corresponds to material ejected at late times. Even if the seed nuclei are as heavy as  $\bar{A} = 120$ , the early episode of mass ejection will have its peak below the Lanthanides.

## 4 SUMMARY AND DISCUSSION

We have explored the effects of a hypermassive neutron star (HMNS) on the long-term evolution of remnant accretion disks formed in neutron star binary mergers. Our main results can be summarized as follows:

1. – A long-lived HMNS results in the ejection of a significant fraction of the disk over a timescale of  $\sim 1$  s. The amount of mass increases monotonically with HMNS lifetime. This enhanced mass loss, up to a factor  $\gtrsim 10$  relative to prompt BH formation, results from enhanced neutrino heating and a reflecting inner boundary.
2. – The composition of the ejecta is latitude-dependent. Material within an angle  $\sim 30^\circ$  of the polar axis is strongly irradiated by the HMNS. If the neutrino and antineutrino luminosities are similar (as we have assumed), the electron fraction of the outflow can be raised to values where Lanthanides are not produced.
3. – Material ejected equatorially is still expected to produce a strong r-process, similar to that of the dynamical ejecta, although the detailed composition is dependent on the HMNS lifetime.

The criterion on the neutron star lifetime to appreciably change the electron fraction of the polar outflow ( $t_{\text{ns}} \gtrsim 100$  ms) is obviously satisfied if the remnant mass is below the maximum mass of a cold neutron star, or if the remnant is stabilized by solid body rotation (a *supramassive* NS), the latter of which is only removed on much longer timescales via e.g. magnetic dipole spin-down. The  $t_{\text{ns}} \gtrsim 100$  ms condition is also likely satisfied if the HMNS is supported by thermal pressure (e.g. Bauswein et al. 2010; Paschalidis et al. 2012; Kaplan et al. 2013), as the latter is removed on the HMNS cooling timescale, which is typically on the order of  $\sim$  seconds (e.g. Pons et al. 1999). On the other hand, support via differential rotation may not last 100 ms, as the latter can be efficiently removed by magnetic stresses: the growth rate of the MRI is as short as milliseconds.

Whether a blue bump indeed develops out of the polar outflow in the case of a long-lived HMNS will depend on a number of factors. First, enough mass needs to be ejected so that the contribution to the lightcurve from radioactive decay becomes detectable. Second, the level of irradiation and the ratio between neutrino and antineutrino luminosities must be sufficient to raise  $Y_e$  to values close to 0.5. The results of Dessart et al. (2009), who employ a much more realistic treatment of neutrino transport but did not include the viscous evolution, seem to align with our findings. The ejecta mass required for a detectable signal depends on how much radioactive heating is supplied by synthesized

<sup>3</sup> We make equations (4)-(5) continuous around  $\bar{Y}_e$  by linear interpolation using the replacements  $(1 - 2Y_e) \rightarrow (1 - Y_e/\bar{Y}_e)$  and  $(\bar{A} - 2\bar{Z}) \rightarrow \bar{A}(1 - \bar{Y}_e/\bar{Y}_e)$  in equations (4) and (5), respectively, when  $Y_e$  is within 10% of  $\bar{Y}_e$ .

<sup>4</sup> Requiring that the argument of the exponential is of order unity would require entropies  $s \gtrsim 600k_{\text{B}}$  per baryon.



elements, which are lighter than the  $r$ -process nuclei but potentially more neutron-rich than  $^{56}\text{Ni}$  (e.g. Grossman et al. 2013).

A final requirement for producing a blue bump is that the high  $Y_e$  (Lanthanide-free) material must be ejected first, such that it resides *exterior* to any lower  $Y_e$  (Lanthanide-rich) material that would otherwise block its emission. Our calculations support this requirement as well: after rising quickly,  $Y_e$  of the polar ejecta decreases monotonically with time (Fig. 5). This is a direct consequence of the decrease in the neutrino luminosities with time, either smoothly as in the case of a long-lived HMNS (Appendix A), or via a sudden drop when a BH forms. Future work will explore in more detail the observational consequences of this bimodal outflow.

Figure 6 provides estimates of various quantities relevant to the kilonova emission derived from our calculations, as a function of the HMNS collapse time. The peak luminosities and time to peak are estimated using the ‘Arnett rule’ for a kilonova (e.g., Li & Paczyński 1998; Metzger et al. 2010),

$$L_{\text{peak}} \simeq 4.3 \times 10^{41} \text{ erg s}^{-1} \left( \frac{f}{3 \times 10^{-6}} \right) \left( \frac{v_r}{0.1c} \right)^{1/2} \times \left( \frac{M_{\text{ej}}}{0.01 M_{\odot}} \right)^{1/2} \kappa^{-1/2} \quad (7)$$

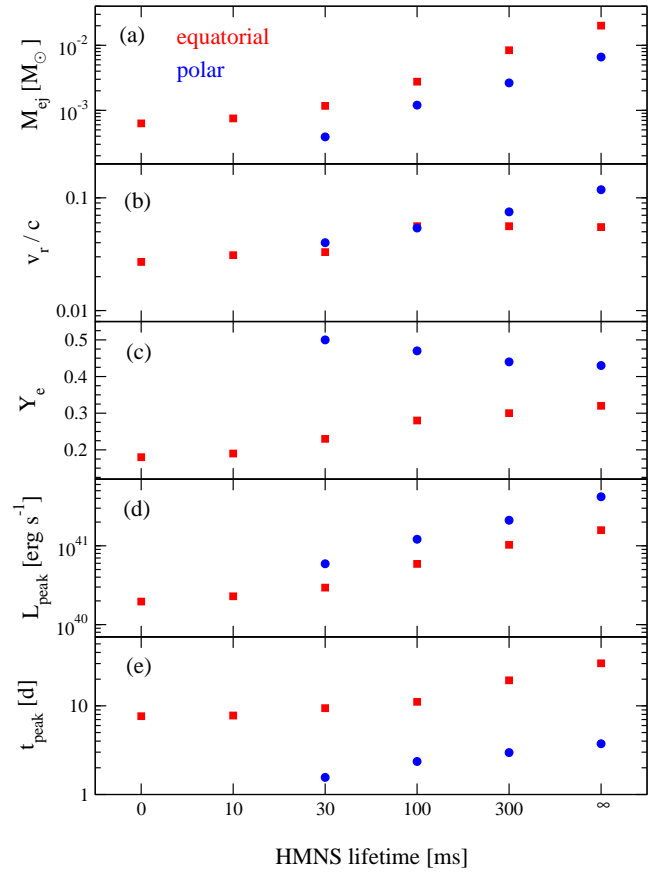
$$t_{\text{peak}} \simeq 1.4 \text{ d} \left( \frac{v_r}{0.1c} \right)^{-1/2} \left( \frac{M_{\text{ej}}}{0.01 M_{\odot}} \right)^{1/2} \kappa^{1/2}, \quad (8)$$

where  $f$  is a factor quantifying radioactive energy deposition, and  $\kappa$  is the opacity of the material in units of  $\text{cm}^2 \text{ g}^{-1}$ . In making Figure 6, we have used  $\kappa = 10 \text{ cm}^2 \text{ g}^{-1}$  for the equatorial material (Lanthanide-dominated), and  $\kappa = 1 \text{ cm}^2 \text{ g}^{-1}$  for the polar material (Lanthanide-free), see Kasen et al. (2013). The velocity is the mass-flux weighted value at  $r = 10^9 \text{ cm}$  considering only unbound material.

Most of the resulting kilonova properties are a monotonic function of the HMNS lifetime. Peak luminosities increase from  $\sim 10^{40}$  to  $\sim 10^{41} \text{ erg s}^{-1}$  for the red component, while the blue component, when present, is brighter by a factor  $\sim 2$ . Similarly, the peak times for the red component range from about a week to a month, increasing with longer HMNS lifetime due to the larger ejected mass. The blue component can last from a few- to several days. While the blue component is faster with longer HMNS lifetime, the average velocity of the red component saturates at  $\sim 0.05c$ .

The large amount of ejecta mass found by our calculations suggests that outflows from the disk could easily overwhelm that from the dynamical ejecta. One implication of this result relates to the Galactic production of  $r$ -process elements. FM13 estimated that ejection of  $\sim 10\%$  of the disk mass would contribute with  $\sim 20\%$  of the production rate of elements with  $A \gtrsim 130$  assuming reasonable values for the disk mass and neutron star merger rate. The increase in the ejected fraction of the disk by a factor of several relative to the prompt BH case implies that disks with long-lived HMNS could become a dominant contribution to the galactic  $r$ -process element production (Freiburghaus et al. 1999; Rosswog et al. 2013; Piran et al. 2014). This is the case even when the polar outflow is Lanthanide free (Fig 4).

The large ejecta masses we infer for a moderately long-lived HMNS ( $t_{\text{ns}} = 100 \text{ ms}$ ) may also help alleviate the ten-



**Figure 6.** Kilonova properties as a function of HMNS lifetime. Shown are the ejecta mass in unbound material (a), the mass-flux weighted velocity at  $10^9 \text{ cm}$  (b), the mass-flux weighted electron fraction (c), the peak luminosity (d) (eq. [7]), and the time to peak (e) (eq. [8]). Red squares and blue dots indicate equatorial and polar material. The opacity is assumed to be Lanthanide-dominated for the former and Lanthanide-free for the latter. Note that these numbers underestimate the ‘red’ component since we have not included contributions from the dynamical ejecta.

sion between large ejecta mass  $\sim 3 \times 10^{-2} M_{\odot}$  required to fit the NIR excess observed following GRB 130603B with current models (Tanaka & Hotokezaka 2013; Piran et al. 2014), without the need to invoke less likely scenarios such as the merger of a NS with a low mass BH. If the long-lived HMNS is magnetized, its rotational spin-down could also power the excess X-rays observed following this event (Fong et al. 2014; Metzger & Piro 2013; Fan et al. 2013). Given that early blue emission appears to require the presence of a long-lived HMNS, detection of such a component in future events provides a relatively clean way to rule out a NS-BH merger (although the *absence* of early blue emission would not rule out a NS-NS merger).

Our models include many approximations in order to make the evolution to a time  $\sim 10 \text{ s}$  computationally feasible. In addition, we have focused here on the key differences introduced by the HMNS applied to models with a particular choice of parameters. Much more work remains in order to make reliable predictions for kilonovae emission. In addition to a more extensive exploration of parameter space, models with realistic angular momentum - and neutrino transport will be needed.

## ACKNOWLEDGMENTS

We thank Dan Kasen, Eliot Quataert, Edo Berger, Gabriel Martínez, Sasha Tchekhovskoy, Wen-fai Fong, and Mansi Kasliwal for helpful discussions and/or comments on the manuscript. BDM acknowledges support from Columbia University. RF acknowledges support from the University of California Office of the President, and from NSF grant AST-1206097. The software used in this work was in part developed by the DOE NNSA-ASC OASCR Flash Center at the University of Chicago. This research used resources of the National Energy Research Scientific Computing Center (NERSC), which is supported by the Office of Science of the U.S. Department of Energy under Contract No. DE-AC02-05CH11231. Computations were performed at *Carver* and *Edison*.

## APPENDIX A: NEUTRINO TREATMENT

### A1 Neutrino Leakage Scheme and Boundary Layer

Following Ruffert et al. (1996), we treat neutrino cooling by smoothly interpolating between the direct neutrino loss rate in the optically thin limit, and the diffusive loss rate of neutrinos in chemical equilibrium for large optical depth. We do this because we reach optical depths  $\gtrsim 10$  in the boundary layer for an extended period of time.

The effective rate per baryon of lepton number loss and the rate per unit mass of energy loss are given respectively by

$$\Gamma_i^{\text{eff}} = \frac{\Gamma_i^0}{1 + t_{\text{n-loss},i}/t_{\text{diff},i}} \quad (\text{A1})$$

$$Q_i^{\text{eff}} = \frac{Q_i^0}{1 + t_{\text{e-loss},i}/t_{\text{diff},i}}, \quad (\text{A2})$$

where the subscript  $i$  refers to electron-type neutrinos or antineutrinos,  $t_{\text{diff},i}$  is the diffusion time, and the direct loss timescales satisfy

$$t_{\text{n-loss},i} = Y_i/\Gamma_i^0 \quad (\text{A3})$$

$$t_{\text{e-loss},i} = e_i/Q_i^0, \quad (\text{A4})$$

with  $\Gamma_i^0$  and  $Q_i^0$  the optically-thin lepton number and energy loss rates, respectively,  $Y_i$  the neutrino number per baryon in chemical equilibrium, and  $e_i$  the specific neutrino energy in chemical equilibrium (Ruffert et al. 1996).

The diffusion time for each neutrino species is given by  $t_{\text{diff},i} = (L/c)\tau_i$ , where  $L$  is a characteristic distance,  $c$  is the speed of light, and  $\tau_i$  is the optical depth of species  $i$  over the distance  $L$ . To keep computations economical, we define the optical depth as

$$\tau_i = \kappa_i \min\{r, H_{\perp}, H_{\parallel}\}, \quad (\text{A5})$$

where  $\kappa_i$  is the corresponding charged-current absorption coefficient, and  $H_{\perp}$  and  $H_{\parallel}$  are the vertical and horizontal gas pressure scale heights, respectively, obtained using the appropriate components of the effective gravitational acceleration vector  $\mathbf{g}_{\text{eff}}$  (e.g.,  $H_{\perp} = p/(\rho|\mathbf{g}_{\text{eff}}|\sin\theta)$ ). This prescription approximates the true optical depth to within a factor of a few. The same length scale used in equation (A5) is assigned to  $L$  for computing the diffusion time  $t_{\text{diff},i}$ .

Figure A1 shows the resulting neutrino luminosities from the disk for different models. Given that at late times most of the emission arises from a thin boundary layer, we have tested the sensitivity of the resulting luminosity to the spatial resolution. Models rA3p15r1 and rA3p15r2 have inner radial cells twice larger and smaller than the fiducial model tInfA3p15, respectively<sup>5</sup>. The luminosities are very close to each other, with an overall trend of less emission with higher resolution. Given this outcome we consider the neutrino emission to be converged. The thermodynamic quantities of the ejecta differ by a few percent or less between models with different resolution. The ejected mass along the equator is the most sensitive, with differences up to a few percent. Model rA3p15r2 is terminated early due to the higher computational cost; comparing the mass outflow rate with that of tInfA3p15 at the same time yields differences up to  $\sim 10\%$  in equatorial ejecta mass.

The effect of the HMNS irradiation on the disk emissivity can be assessed by comparing the luminosity from model pA3p15xs with that of the fiducial model tInfA3p15. When the HMNS emits, the disk luminosity increases relative to no HMNS irradiation over the thermal time of the disk ( $\sim 30$  ms at  $r \sim 100$  km). The output luminosity exceeds the imposed HMNS flux due to the contribution from accretion.

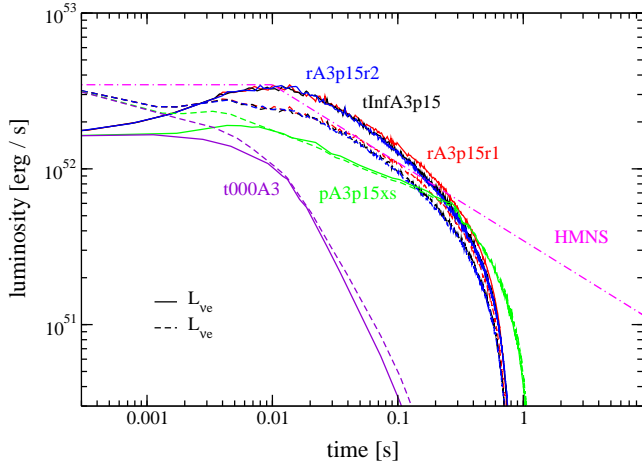
The effect of the reflecting boundary can be quantified by comparing the luminosity of model pA3p15xs with that of model t00A3. The increased radiative efficiency results in higher emission by about an order of magnitude at time  $\sim 0.1$ s, explaining why both neutrino energy deposition and composition changes are more prominent when the HMNS is present.

### A2 Self-Irradiation with a Boundary Layer

The spatial emission properties of the disk are slightly more complicated when a HMNS sits at the center than when a BH is the central object. Most of the emission comes from two hot spots located above and below the optically-thick midplane (Fig. 2c). Thus our prescription for *disk self-irradiation* requires modifications relative to FM13. Since the position of these hot spots evolves in both radius and polar angle, we compute the neutrino temperature as an emissivity-weighted average over the entire computational domain instead of a mass-weighted average. This way we avoid biasing this temperature towards the optically thick midplane, which does not emit much but has the highest temperature. We ignore differences between neutrino and antineutrino temperature from the disk.

Given that the emission is equatorially symmetric, we use the same angular distribution as in FM13, with the emission radius  $R_{\text{em}}$  computed in the same way. However, the attenuation of the radiation flux from the disk is modified in two ways. First, we multiply by an overall factor of  $1/2$  to account for the fact that the midplane acts as a radiation insulator. Second, we use an attenuation optical depth consistent with that employed in the leakage calculation.

<sup>5</sup> Since the radial grid is ratioed, this means that the increased resolution in inner regions is compensated by lower resolution outwards for a fixed number of cells.



**Figure A1.** Neutrino and antineutrino luminosities from the disk as a function of time (solid and dashed lines, respectively). Shown are three models with different radial resolutions near the inner boundary (rA3p15r1, tInfA3p15, rA3p15r2), a model with irradiation from the HMNS suppressed, keeping the reflecting boundary and disk self-irradiation (pA3p15xs), and a model with a BH at the center (t000A3). The imposed HMNS luminosity (equal for neutrinos and antineutrinos) is shown by the dot-dashed magenta line.

That is, the radiation flux from the disk seen by a given point in the computational domain is multiplied by a factor  $\exp(-\tau_{\text{irr},i})$ , with  $\tau_{\text{irr}}$  being the maximum of the values obtained by evaluating equation (A5) at the location of the emission hot spot and at the absorption point. This prescription is good at capturing large density contrasts, but fails to account for shadowing at large radii along the equator.

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